

Photogalvanic current in a double quantum well

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We study the in-plane stationary current caused by phototransitions between the states of a double quantum well. The electric polarization of light has both vertical and in-plane components. The stationary current originates from the periodic vibration of electrons between two non-equivalent quantum wells caused by the normal component of the alternating electric field with simultaneous in-plane acceleration/deceleration by the in-plane component of electric field. The quantum mechanism of the stationary current is conditioned by in-plane transition asymmetry which appears due to the indirect phototransitions with the participation of impurity scattering. The photocurrent has a resonant character corresponding to the equality of the photon energy to the distance between subbands. It is found that the current appears as a response to the linear-polarized light.

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I. INTRODUCTION

Since the first studies on the photogalvanic effect (PGE) at the end of the 70th, a wide literature devoted to this subject has appeared [1–3], see also reviews [4–9]. The activity in this field continues up to now (see, for example, [10–18]). There are different variants of PGE in confined systems: the stationary in-plane photocurrent in classical [19] and quantum [20, 21] films, and the current along solid-state surface [22–24]. This photocurrent exists even if crystal asymmetry is negligible, but the quantum well is oriented (directions across the well are not equivalent). The current along the surface occurs if the electric field of the light has both in- and out-plane components.

The phenomenology of PGE is determined by the relation for the current density

$$\mathbf{j} = \alpha_s((\mathbf{E} - \mathbf{n}(\mathbf{nE}))(\mathbf{nE}^* + \mathbf{c.c}) + \mathbf{i}\alpha_a[\mathbf{n}[\mathbf{E}\mathbf{E}^*]]), \quad (1)$$

where \mathbf{n} is the normal to the quantum well. Real constants α_s and α_a describe linear and circular photogalvanic effects, correspondingly. The origin of this current can be understood if to consider the out-of-plane component of electric field as modulating the quantum well conductivity with a simultaneous driving of electrons by the field in-plane.

In a quantum well the vertical component of the electric field of light can cause the transitions between different quantum subbands. In the presence of scattering this gives birth to the effective pumping of the in-plane momentum to the electronic system. The light plays the role of the energy and non-equilibrium source, while the scatterers produce electrons in-plane acceleration. The situation is, in a certain sense, similar to the motion of a car where the friction forces the car to move.

The purpose of the present article is to study the mechanism of PGE in a double quantum well. This system looks perspective because the structure of the levels of a double quantum well permits easy tuning of the distance between subbands to the frequency of the external field.

The effect under consideration is illustrated in Fig.1. We consider intersubband transitions of electrons in a system with the quadratic energy spectrum. An electron goes between two states $\epsilon_n(\mathbf{p})$ and $\epsilon_{n'}(\mathbf{p}')$ due to the simultaneous action of electric field and scattering. These states originate from mixing the states of different individual quantum wells. The in-plane current appears due to the change of electron in-plane momentum. To "memorize" electric field in- and out-plane components, the transition probability should contain their product. For non-conservation of the electron momentum the scattering should be taken into account. This transition probability arises in the second order of the perturbation theory. The amplitude of transitions has a resonance on an intermediate state. The subbands of the quantum well are equidistant, that gives rise to the absence of the resonance smearing due to the difference in electron momenta. The result of excitation is the pumping of the momentum to the electron subsystem and the in-plane current. The paper is organized as follows. First, we will discuss a simple classical model of the effect based on a parabolic well. Then, we will find the transition probability in a classical electric field. After that the current will be found using many-band kinetic equation.

II. SIMPLE CLASSICAL MODEL

To explain the physical origin of the effect we consider a simple classical model instead of a 2D system: an electron in an oscillatory well in z -direction with confining potential $m\omega_0^2 z^2/2$ affected by the alternating electric field with x and z components $\mathbf{E}(t) = \text{Re}(\mathbf{E}e^{-i\omega t})$. The classical Newton equation for an electron reads

$$\ddot{\mathbf{r}} + \gamma \dot{\mathbf{r}} = e\mathbf{E}/m, \quad (2)$$

where we introduced the liquid friction coefficient $\gamma = \gamma_0 + \gamma_1 z$. The dependence of the friction on z takes into account the assumed weak asymmetry ($\gamma_0 \gg \gamma_1 z$) of the well in z -direction.

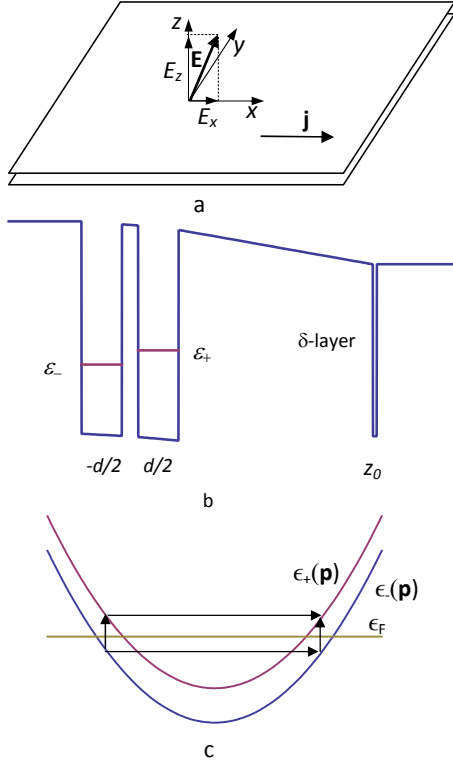


FIG. 1: (Color online) a) The sketch of the proposal experimental setup. The electric field of light $\mathbf{E}(t)$ is tilted in (x,z) plane. The stationary current is directed along the x -axis. b) The sketch of the band structure. Quantum wells are centered in planes $z = \pm d/2$. The carriers are provided by the δ -layer of donors in plane z_0 . c) The transition amplitude includes vertical transition caused by light between \pm subbands and impurity scattering which does not conserve the in-plane momentum.

The forced solution of the Newton equation is found by expanding in powers of γ_1 :

$$\mathbf{r} = \mathbf{r}_0 + \mathbf{r}_1 + \dots,$$

$$z_0 = \text{Re} \frac{eE_z}{m(-\omega^2 + \omega_0^2 - i\gamma_0\omega)} e^{-i\omega t},$$

$$x_0 = \text{Re} \frac{eE_x}{m(-\omega^2 - i\gamma_0\omega)} e^{-i\omega t},$$

$$\overline{(\dot{x}_1)} = \frac{\gamma_1 \omega e^2}{2\gamma_0 m^2} \text{Im} \frac{E_x^* E_z}{(\omega^2 - i\gamma_0\omega)(\omega^2 - \omega_0^2 + i\gamma_0\omega)}, \quad \overline{(\dot{z}_1)} = 0. \quad (3)$$

Here $\overline{(\dots)}$ denotes the time averaging.

Let damping γ_0 be also small. Then the mean velocity has a resonance at $\omega = \omega_0$. The frequency behavior near this point depends on the kind of electromagnetic field polarization: delta-like peak for the linear polarization and antisymmetric Fano-like resonance $\propto 1/(\omega - \omega_0)$ for

circular polarization. The origin of this behavior is explained by the character of the electron motion in the zero approximation. Indeed, if $\gamma_1 = 0$, for linear polarization, the electron rotates in the exact resonance and vibrates along a straight line out of resonance. For circular polarization the behavior is opposite.

Liquid friction force $-\gamma \dot{\mathbf{r}}$ does not affect the direction of vibrating motion; therefore it does not produce a drift. At the same time, due to γ_1 , a rotating particle differently brakes at the opposite (upper and lower) sides of the circle that produces a translational displacement, and as a result, the mean drift. In the case of circular-polarized light, the direction of the motion depends on the sign of polarization and the sign of resonance detuning. The value of the drift velocity near resonance does not depend on the friction strength, but it depends on ratio γ_1/γ .

Fig.2 illustrates this reasoning by the exact solution of the Newton equation. The photogalvanic effect in this

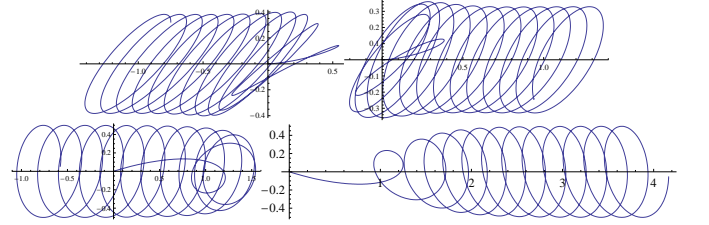


FIG. 2: (Color online) Solution of the Newton equation for linear (c,d) and circular (a,b) polarized electric field for initial conditions $\mathbf{r}(0) = 0$, $\dot{\mathbf{r}}(0) = 0$ and parameters $\gamma_0 = 0.2$, $\gamma_1 = 0.02$, $eE_x/m = 0.3$, $eE_z/m = 0.1$, $\omega_0 = 1$; for plot a) $\omega = 0.9$, for plot b) $\omega = 1.1$, for plots a) and b) $\omega = 1$. For circular polarization the sign of detuning determines the direction of the steady-state drift.

model has a purely classical nature. In particular, the circular PGE does not need the spin pumping as in spin-related circular PGE. At the same time, the classical and quantum photogalvanic effects have different properties. The photogalvanic effect on intersubband transitions of double quantum well Considered below has a resonant character like the classical PGE discussed here. The difference is the absence of circular photogalvanic effect for transitions in a double quantum well.

III. TRANSITIONS BETWEEN SUBBANDS OF DOUBLE WELL

We study electrons with a parabolic isotropic energy spectrum in a double quantum well (see Fig.1). The amplitude of transition between wells is weak, but comparable to the separation of energies of individual wells. The states with the in-plane electron momentum \mathbf{p} and the subband number $n = \pm |n, \mathbf{p}\rangle = \chi_n(z) \exp(i\mathbf{p}\boldsymbol{\rho})/\sqrt{S}$, (S is the system area, we set $\hbar = 1$ throughout this section besides the final expression) have energies $\epsilon_{n,\mathbf{p}} = p^2/2m + \epsilon_n$. In this case, the subbands are parallel,

$\epsilon_{+, \mathbf{p}} - \epsilon_{-, \mathbf{p}} \equiv \epsilon_+ - \epsilon_-$. This circumstance plays an important role in the further consideration, providing the resonance of optical frequency with a distance between subbands for electrons with arbitrary momenta. The overlapping of wave functions $\chi_n(z)$ is supposed to be weak and intersubband distance $\epsilon_+ - \epsilon_- = \Delta$ ($\Delta > 0$) is small as compared to the Fermi energy. The scatterers (donors) are distributed in a delta-layer at $z = z_0 > 0$. The well widths and the distance d between them are assumed to be small as compared to z_0 .

Assuming that the mean free time is large as compared to the distance between the levels of quantum wells (and also the Fermi energy) one can treat n and \mathbf{p} as good quantum numbers and describe the problem within the kinetic equation for distribution functions $f_{n, \mathbf{p}}$. In such an equation, external classical alternating electric field $\mathbf{E}(t) = \text{Re}(\mathbf{E}_0 e^{-i\omega t})$ causes the transition between unperturbed states and determines the generation term in the kinetic equation. The interaction with charged impurities provides the mechanism of electron scattering. The kinetic equation reads

$$\sum_{n', \mathbf{p}'} W_{n, \mathbf{p}; n', \mathbf{p}'}^{imp} (f_{n', \mathbf{p}'}^{(1)} - f_{n, \mathbf{p}}^{(1)}) + G_{n, \mathbf{p}} = 0, \quad (4)$$

where the generation $G_{n, \mathbf{p}}$ is given by

$$G_{n, \mathbf{p}} = \sum_{n', \mathbf{p}'} W_{n, \mathbf{p}; n', \mathbf{p}'}^{ph} (f_{n', \mathbf{p}'}^{(0)} - f_{n, \mathbf{p}}^{(0)}). \quad (5)$$

Here $W_{n, \mathbf{p}; n', \mathbf{p}'}^{imp}$ is the impurity transition probability, $W_{n, \mathbf{p}; n', \mathbf{p}'}^{ph}$ is the transition probability due to the combined action of electromagnetic field and impurities, $f_{n, \mathbf{p}}^{(0)}$ is the equilibrium distribution function and $f_{n, \mathbf{p}}^{(1)}$ is the first correction to the distribution function in the external electromagnetic field. Using the classical kinetic equation means neglecting the off-diagonal elements of the density matrix that is valid if the collision broadening of subbands is less than the distance between them. The perturbation includes the Hamiltonian of the interaction with electromagnetic field \hat{H}^{ph} and the potential energy of the electron interacting with impurities \hat{V} . The first is

$$\hat{H}^{ph} = \frac{e}{c} \text{Re}(\mathbf{A} e^{-i\omega t}) \hat{\mathbf{v}} \equiv \frac{1}{2} (\hat{U} e^{-i\omega t} + h.c.), \quad (6)$$

where $\text{Re}(\mathbf{A} e^{-i\omega t})$ is the vector potential of electromagnetic field with frequency ω , $\hat{\mathbf{v}} = (\hat{\mathbf{v}}^{\parallel}, \hat{\mathbf{v}}^z)$ is the velocity operator. The complex amplitude of electric field is $\mathbf{E} = i\omega \mathbf{A}/c$. Thus, the operator $\hat{U} = e(\mathbf{E} \hat{\mathbf{v}})/i\omega$. Note that we suppose the electric field to be homogeneous. The diagonal elements of in-plane components of the the velocity operator $\mathbf{v}_{n, \mathbf{p}; n', \mathbf{p}'}^{\parallel} = \mathbf{v}_{\mathbf{p}} \delta_{n\mathbf{p}, n'\mathbf{p}'}$, $\mathbf{v}_{\mathbf{p}} = \partial_{\mathbf{p}} \epsilon_{n, \mathbf{p}} = \mathbf{p}/m$. The normal component has matrix elements $v_{n, \mathbf{p}; n', \mathbf{p}'}^z = v_{n, n'}^z \delta_{\mathbf{p}, \mathbf{p}'}$. The impurity potential reads

$$V(\mathbf{r}) = \sum_i u(\mathbf{r} - \mathbf{r}_i), \quad (7)$$

where the sum runs over all the impurities situated in points \mathbf{r}_i with individual potentials $u(\mathbf{r} - \mathbf{r}_i)$.

The appearance of the photogalvanic current requires non-conservation of the in-plane momentum in the electron excitation process. Hence, the phototransitions should include the participation of the "third body". In our case the impurities play the role of this agent. The excitation probability including the impurity scattering is determined by the second-order transition amplitude. The needed term arises from the interference of amplitudes caused by the E_z and in-plane components of the electric field. The draft of the transitions is depicted in Fig.1.

In the second order of the interaction, the transition probability is

$$\begin{aligned} W_{n, \mathbf{p}; n', \mathbf{p}'}^{ph} = & \frac{\pi}{2} \left\langle \left| \sum_{n_1} \left(\frac{V_{n, \mathbf{p}; n_1, \mathbf{p}'} U_{n_1, \mathbf{p}; n', \mathbf{p}'}^+}{\eta + i(\epsilon_{n_1, n'} + \omega)} + \frac{U_{n, \mathbf{p}; n_1, \mathbf{p}}^+ V_{n_1, \mathbf{p}; n', \mathbf{p}'} }{\eta + i(\epsilon_{n_1, n} - \omega)} \right) \right|^2 \right\rangle \times \\ & \delta(\epsilon_{n, \mathbf{p}} - \epsilon_{n', \mathbf{p}'} + \omega) + \\ & \frac{\pi}{2} \left\langle \left| \sum_{n_1} \left(\frac{V_{n, \mathbf{p}; n_1, \mathbf{p}'} U_{n_1, \mathbf{p}; n', \mathbf{p}'} }{\eta + i(\epsilon_{n_1, n'} - \omega)} + \frac{U_{n, \mathbf{p}; n_1, \mathbf{p}} V_{n_1, \mathbf{p}; n', \mathbf{p}'} }{\eta + i(\epsilon_{n_1, n} + \omega)} \right) \right|^2 \right\rangle \times \\ & \delta(\epsilon_{n, \mathbf{p}} - \epsilon_{n', \mathbf{p}'} - \omega); \quad (\eta = +0). \end{aligned} \quad (8)$$

Here $\epsilon_{n_1, n} \equiv \epsilon_{n_1} - \epsilon_n$; angular brackets denote the average over impurities configuration. Using relations $U_{n, \mathbf{p}; n', \mathbf{p}'}^+ = (U_{n', \mathbf{p}'; n, \mathbf{p}})^*$, $V_{n, \mathbf{p}; n', \mathbf{p}'} = (V_{n', \mathbf{p}'; n, \mathbf{p}})^*$ it is easy to prove that $W_{n, \mathbf{p}; n', \mathbf{p}'}^{ph} = W_{n', \mathbf{p}'; n, \mathbf{p}}^{ph}$.

The denominators in Eq.(8) have their resonance with the field frequency independently from the electron momentum. At the same time, the resonance in the final state is absent due to non-conservation of the in-plane momentum.

Eq.(8) can be rewritten in the form ($\mathbf{E} = (\mathbf{E}_{\parallel}, E_z)$):

$$\begin{aligned} W_{n, \mathbf{p}; n', \mathbf{p}'}^{ph} = & \frac{\pi e^2}{2\omega^2} \left\langle \left| \sum_{n_1} \left(V_{n, \mathbf{p}; n_1, \mathbf{p}'} \left(\frac{\mathbf{v}_{\mathbf{p}'} \mathbf{E}_{\parallel}^* \delta_{n_1, n'}}{i\omega} + \frac{v_{n_1, n'}^z E_z^*}{\eta + i(\epsilon_{n_1, n'} + \omega)} \right) + \right. \right. \\ & \left. \left. + \left(\frac{\mathbf{v}_{\mathbf{p}} \mathbf{E}_{\parallel}^* \delta_{n, n_1}}{-i\omega} + \frac{v_{n, n_1}^z E_z^*}{\eta + i(\epsilon_{n_1, n} - \omega)} \right) V_{n_1, \mathbf{p}; n', \mathbf{p}'} \right) \right|^2 \right\rangle \times \\ & \delta(\epsilon_{n, \mathbf{p}} - \epsilon_{n', \mathbf{p}'} + \omega) + \\ & \frac{\pi}{2} \left\langle \left| \sum_{n_1} \left(V_{n, \mathbf{p}; n_1, \mathbf{p}'} \left(\frac{\mathbf{v}_{\mathbf{p}'} \mathbf{E}_{\parallel} \delta_{n_1, n'}}{-i\omega} + \frac{v_{n_1, n'}^z E_z}{\eta + i(\epsilon_{n_1, n'} - \omega)} \right) + \right. \right. \\ & \left. \left. + \left(\frac{\mathbf{v}_{\mathbf{p}} \mathbf{E}_{\parallel} \delta_{n, n_1}}{i\omega} + \frac{v_{n, n_1}^z E_z}{\eta + i(\epsilon_{n_1, n} + \omega)} \right) V_{n_1, \mathbf{p}; n', \mathbf{p}'} \right) \right|^2 \right\rangle \times \\ & \delta(\epsilon_{n, \mathbf{p}} - \epsilon_{n', \mathbf{p}'} - \omega). \end{aligned} \quad (9)$$

It is evident that the contribution to photogalvanic effect is given by not the total transition probability W^{ph} but only its odd in \mathbf{p}, \mathbf{p}' part. For this part, we have the

following expression:

$$\begin{aligned} \tilde{W}_{n,\mathbf{p};n',\mathbf{p}'}^{ph} = & \frac{\pi e^2}{\omega^3} \left\{ \left\langle \text{Re} \left[\sum_{n_1} \left(V_{n,\mathbf{p};n',\mathbf{p}'}(\mathbf{p}' - \mathbf{p}) \mathbf{E}_{\parallel}^* \times \right. \right. \right. \right. \\ & \left. \left. \left(\frac{V_{n,\mathbf{p};n_1,\mathbf{p}'}^* v_{n',n_1}^z E_z}{i\eta + (\varepsilon_{n_1,n'} + \omega)} + \frac{v_{n_1,n}^z E_z V_{n_1,\mathbf{p};n',\mathbf{p}'}^*}{i\eta + (\varepsilon_{n_1,n} - \omega)} \right) \right] \right\rangle \times \\ & \delta(\varepsilon_{n,\mathbf{p}} - \varepsilon_{n',\mathbf{p}'} + \omega) + \left\langle \text{Re} \left[\sum_{n_1} \left(V_{n,\mathbf{p};n',\mathbf{p}'}(\mathbf{p} - \mathbf{p}') \mathbf{E}_{\parallel} \times \right. \right. \right. \\ & \left. \left. \left(\frac{V_{n,\mathbf{p};n_1,\mathbf{p}'}^* v_{n',n_1}^z E_z^*}{i\eta + (\varepsilon_{n_1,n'} - \omega)} + \frac{v_{n_1,n}^z E_z^* V_{n_1,\mathbf{p};n',\mathbf{p}'}^*}{i\eta + (\varepsilon_{n_1,n} + \omega)} \right) \right] \right\rangle \times \\ & \left. \delta(\varepsilon_{n,\mathbf{p}} - \varepsilon_{n',\mathbf{p}'} - \omega) \right\}. \end{aligned} \quad (10)$$

Kinetic equation Eq.(4) can be transformed to

$$\frac{1}{\tau_n(p)} f_{n,\mathbf{p}}^{(1)} - \frac{1}{\tau_{n,-n}(p)} f_{-n,\mathbf{p}}^{(1)} = G_{n,\mathbf{p}}, \quad (11)$$

where $\tau_n(p)$ is the intra-subband transport relaxation time and $\tau_{n,-n}(p)$ is the time of transition from the state (n, \mathbf{p}) to all states of the subband $(-n)$. These values are determined by

$$\begin{aligned} \frac{1}{\tau_n(p)} &= 2\pi \sum_{\mathbf{p}'} \left[\left\langle |V_{n,\mathbf{p};n,\mathbf{p}'}|^2 \right\rangle \delta(\varepsilon_{n,\mathbf{p}} - \varepsilon_{n,\mathbf{p}'})(1 - \frac{\mathbf{p}\mathbf{p}'}{p^2}) \right. \\ &\quad \left. + \left\langle |V_{n,\mathbf{p};-n,\mathbf{p}'}|^2 \right\rangle \delta(\varepsilon_{n,\mathbf{p}} - \varepsilon_{-n,\mathbf{p}'} \right]; \\ \frac{1}{\tau_{n,-n}(p)} &= 2\pi \sum_{\mathbf{p}'} \left\langle |V_{n,\mathbf{p};-n,\mathbf{p}'}|^2 \right\rangle \delta(\varepsilon_{n,\mathbf{p}} - \varepsilon_{-n,\mathbf{p}'} \frac{\mathbf{p}\mathbf{p}'}{p^2}) \end{aligned} \quad (12)$$

Solving Eq.(11) we find (argument \mathbf{p} is omitted):

$$f_n^{(1)} = \left(G_n \tau_n + G_{-n} \frac{\tau_+ \tau_-}{\tau_{n,-n}} \right) \left(1 - \frac{\tau_+ \tau_-}{\tau_{n,-n}} \right)^{-1}. \quad (13)$$

The expressions for $\tau_n(p)$, $\tau_{n,-n}(p)$ and $\tilde{W}_{n,\mathbf{p};n',\mathbf{p}'}^{ph}$ contain correlators of the form $\langle V_{n,\mathbf{p};n',\mathbf{p}'} V_{m,\mathbf{p};m',\mathbf{p}'}^* \rangle$. In the case of impurities situated in layer $z = z_0$ ($\mathbf{r}_i = (\boldsymbol{\rho}_i, z_0)$) the function $V(\mathbf{r})$ reads

$$V(\mathbf{r}) = \sum_{\mathbf{q},i} u_{\mathbf{q}} e^{-q|z-z_0|} \exp(-i\mathbf{q}(\boldsymbol{\rho} - \boldsymbol{\rho}_i)), \quad (14)$$

where $u_{\mathbf{q}}$ is the 2D Fourier component of the impurity center potential. For example, for unscreened Coulomb center $u_{\mathbf{q}} = 2\pi e^2 / \kappa q S$ (κ is the background dielectric constant). Correlators are given by

$$\begin{aligned} \langle V_{n,\mathbf{p};n',\mathbf{p}'} V_{m,\mathbf{p};m',\mathbf{p}'}^* \rangle &= n_s S \int dz dz' |u_{\mathbf{p}-\mathbf{p}'}|^2 e^{-q(2z_0-z-z')} \\ &\quad \times \chi_n(z) \chi_{n'}(z) \chi_m(z') \chi_{m'}(z'). \end{aligned} \quad (15)$$

Here n_s is the areal density of scatterers. We suppose that the electron wavelength is larger than d . In this approximation one can find from Eq.(15):

$$\langle V_{n,\mathbf{p};n',\mathbf{p}'} V_{m,\mathbf{p};m',\mathbf{p}'}^* \rangle = n_s S |u_{\mathbf{p}-\mathbf{p}'}|^2 e^{-2qz_0} \times \left[\delta_{n,n'} \delta_{m,m'} + q(z_{n,n'} \delta_{m,m'} + z_{m,m'} \delta_{n,n'}) \right]. \quad (16)$$

Matrix elements $z_{nn'}$ should be estimated for specific wave functions. For simplicity, we will use the wave functions of two delta-functional wells in the tight-binding approximation. The seed states with energies $\varepsilon_0 \pm \Delta_0/2$ can be written as

$$\chi_{1,2} = \sqrt{\kappa} e^{-\kappa|z \mp d/2|}. \quad (17)$$

In basis (17) $\chi_+ = (1, \beta) / \sqrt{1 + \beta^2}$, $\chi_- = (\beta, -1) / \sqrt{1 + \beta^2}$, where β is the mixing amplitude. The corresponding states energies are $\varepsilon_{\pm} = \varepsilon_0 \pm \Delta/2$, $\Delta = \sqrt{\Delta_0^2 + 4t_0^2}$, where $t_0 \sim \varepsilon_0 e^{-\kappa d}$ is a hopping amplitude between wells. For quantity β we have $\beta = 2t_0 / (\Delta + \Delta_0)$. The matrix elements of z are $z_{++} = -z_{--} = d(1 - \beta^2) / (2(1 + \beta^2))$, $z_{+-} = \beta d / (1 + \beta^2)$.

Inserting Eq.(16) in Eq.(12) we get the expressions for $\tau_+ \approx \tau_- = \tau$ and a small difference $1/\tau_- - 1/\tau_+$:

$$\begin{aligned} \frac{1}{\tau} &= mn_s \int \frac{d\mathbf{q}}{2\pi} |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0} \delta(q^2 + 2\mathbf{p}\mathbf{q}) \frac{q^2}{p^2} \\ \frac{1}{\tau_-} - \frac{1}{\tau_+} &= m(z_{++} - z_{--}) n_s \times \\ &\int \frac{d\mathbf{q}}{\pi} |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0} \delta(q^2 + 2\mathbf{p}\mathbf{q}) \frac{q^3}{p^2}, \end{aligned} \quad (18)$$

where $\tilde{u}_{\mathbf{q}} = S u_{\mathbf{q}}$. From Eq.(16) it is seen that $\tau_{n,-n} \gg \tau_n$ and, so, Eq.(13) can be simplified

$$f_n^{(1)} = G_n \tau_n. \quad (19)$$

Further we will consider the resonance situation when frequency ω is close to Δ . Smallness Δ , as compared to the Fermi energy $\epsilon_F = p_F^2 / 2m$ (p_F being the Fermi momentum) leads to approximate expressions for $G_+ \approx -G_-$,

$$\begin{aligned} G_+ &= \frac{n_s e^2 \Delta z_{+-}^2}{\pi \omega^2} \frac{\eta}{(\Delta - \omega)^2 + \eta^2} \frac{\partial f_{\mathbf{p}}^{(0)}}{\partial \mu} \times \\ &\int d\mathbf{q} (\mathbf{q} \cdot \text{Re}(\mathbf{E}_{\parallel} E_z^*)) q |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0} \delta(q^2 + 2\mathbf{p}\mathbf{q}). \end{aligned} \quad (20)$$

As a result, for the current of photogalvanic effect, we have

$$\begin{aligned} \mathbf{j} &= -\frac{n_s e^3 \Delta z_{+-}^2}{4\pi^3 m \omega^2} \frac{(\tau_+ - \tau_-) \tau}{(\Delta - \omega)^2 \tau^2 + 1} \text{Re}(\mathbf{E}_{\parallel} E_z^*) \times \\ &\int d\mathbf{p} \frac{\partial f_{\mathbf{p}}^{(0)}}{\partial \mu} \int d\mathbf{q} |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0} q^3 \delta(q^2 + 2\mathbf{p}\mathbf{q}). \end{aligned} \quad (21)$$

Eq. (21) has a resonant character with the resonance at $\omega = \Delta$. This resonance results from the intermediate

state for transition due to the parallelism (equidistance) of subbands. The resonance is smeared due to scattering, e.g., by impurities. To include this smearing, the infinitesimal η was replaced by finite relaxation rate $1/\tau$ which can be estimated from mobility. This leads to the finiteness of the current at the point of resonance.

At temperature $T = 0$ the latter expression is simplified, and we obtain the final result for the required value:

$$\mathbf{j} = -\frac{4e^3(z_{++} - z_{--})\Delta z_{+-}^2 \epsilon_F \tau}{\pi \omega^2 d^2 ((\Delta - \omega)^2 \tau^2 + 1)} \text{Re}(\mathbf{E}_{\parallel} E_z^*) F, \quad (22)$$

where we introduced a dimensionless quantity $F = d^2 \Phi_3^2 \Phi_2^{-2}$,

$$\Phi_s = \int_0^{2p_F} dq q^s |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0} \frac{1}{\sqrt{1 - q^2/4p_F^2}}. \quad (23)$$

In the specific case of $p_F z_0 \gg 1$ Eq. (23) is reduced to

$$\Phi_s = \int_0^{\infty} dq q^s |\tilde{u}_{\mathbf{q}}|^2 e^{-2qz_0}. \quad (24)$$

If the scattering is determined by the charged non-screened impurities $F = d^2/4z_0^2$.

In the model of two δ -like wells Eq.(22) gives For linear polarized wave

$$\mathbf{j} = -\frac{4e^3 d \beta (1 - \beta^2) n}{m(1 + \beta^2)^2 \Delta} \frac{\tau}{(\Delta/\hbar - \omega)^2 \tau^2 + 1} E_0^2 \sin(2\theta) F, \quad (25)$$

where E_0 is the amplitude of the electric field, θ is the angle between the field and the normal to the system.

It should be emphasized that the current contains the linear response only. This distinguishes the quantum double-well result from the simple classical model of the effect considered in the previous section.

Let us estimate the value of the effect. Considering β as a free parameter we can choose $\beta = \sqrt{2}-1$ to maximize the β -dependent factor in Eq.(25) $\beta(1 - \beta^2)/(1 + \beta^2)^2 =$

$1/4$. The optimum for PGE observation corresponds to $\omega = \Delta$ and $\theta = \pi/4$. Choosing the typical values for *GaAs/AlGaAs* double quantum wells $d = 5 \cdot 10^{-7}$ cm, $z_0 = 3 \cdot 10^{-6}$ cm, $\epsilon_F = 20$ meV ($n = 6.2 \cdot 10^{11} \text{ cm}^{-2}$), $\Delta = 0.1$ meV, $\tau = 4 \cdot 10^{-11}$ s, $E_0 = 1$ V/cm we find for this optimal situation $j \approx 3.6 \mu\text{A/cm}$ that is a quite measurable value.

It should be emphasized that the initial and final states in the transition can belong to the different or the same subbands. The resonant behavior results from the resonance on the intermediate state rather than the energy conservation in the final states, because the conservation law for the phototransition with the participation of impurity scattering does not give a fixed frequency for the transition. The sharpness of the resonance is conditioned by the equidistance of the energy bands in a 2D well.

IV. CONCLUSIONS

We found the stationary current along a double-well system affected by the linear-polarized far-infrared wave. The stationary current originates from the periodic vibration of electrons between two non-equivalent quantum wells caused by the normal component of the alternating electric field with synchronic in-plane acceleration/deceleration by the in-plane component of the electric field. The linear photogalvanic effect needs vertical asymmetry of the quantum well. The effect has the peak resonant structure connected with the parallel subbands of the double quantum well. The resonant frequency can be easily tuned by the application of the gate voltage. The optimal range of frequencies is $10^{11} \div 10^{13} \text{ s}^{-1}$. The predicted value of the current is experimentally measurable.

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